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# Radiative Gaugino Masses

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**Abstract:** We investigate the possibility that gauginos are massless at tree level and that the U(1) R-invariance is broken spontaneously by Higgs vevs, like the chiral symmetry of quarks in the standard model, or else explicitly by dimension 2 or 3 SUSY-breaking terms in the low energy effective Lagrangian. Gluino and lightest neutralino masses then depend on only a few parameters. For a SUSY-breaking scale  $\lesssim 400$  GeV, the gluino and lightest neutralino have masses typically in the range  $1/10 \sim 2\frac{1}{2}$  GeV. On the other hand, for a SUSY-breaking scale several TeV or larger, radiative contributions can yield gluino and lightest neutralino masses of  $O(50 - 300)$  GeV and  $O(10 - 30)$  GeV, respectively. As long as the Higgs vev is the only source of R-invariance breaking, or if SUSY breaking only appears in dimension 2 terms in the effective Lagrangian, the gluino is generically the lightest SUSY particle, modifying the usual phenomenology in interesting ways.

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# 1 Introduction

There is nowadays an intense effort to understand the nature and structure of the supersymmetry breaking sector in low energy effective theories which are obtained as the pointlike limit of superstrings. The pattern of these soft breaking terms is obviously linked to the mechanism which is chosen in superstrings to originate the breaking of the local supersymmetry. One interesting type of SUSY breaking predicts vanishing gaugino masses at the scale of supergravity breaking. This class of superstring models is often discarded on the phenomenological basis that gaugino masses (in particular the gluino mass) in the low-energy theory would be too small. In this paper we discuss this possibility and we show that scenarios with vanishing tree-level gaugino masses are not so strictly excluded as is commonly believed.

R-invariance is automatically a symmetry of the MSSM Lagrangian before supersymmetry is broken. In superfield form, the F-terms of this Lagrangian have the trilinears which are needed to give ordinary fermions their masses:  $(\hat{Q}\hat{U}^c\hat{H}_u)_{\theta\theta}$  and  $(\hat{Q}\hat{D}^c\hat{H}_d)_{\theta\theta}$  and the analogs for the leptons. In addition, the term  $\mu(\hat{H}_u\hat{H}_d)_{\theta\theta}$  is needed to break the ew gauge symmetry, at least in the scenario of refs. [1, 2]. Assigning  $R(\theta) = 1$ , an R-charge assignment for the chiral superfields can be found which satisfies the conditions

$$\begin{aligned} R(\hat{H}_u) + R(\hat{H}_d) &= 2 \\ R(\hat{H}_u) + R(\hat{Q}) + R(\hat{U}^c) &= 2 \\ R(\hat{H}_d) + R(\hat{Q}) + R(\hat{D}^c) &= 2 \\ R(\hat{H}_d) + R(\hat{L}) + R(\hat{E}^c) &= 2 \end{aligned} \tag{1}$$

so that the Lagrangian is R-invariant<sup>2</sup>. Some soft-SUSY-breaking terms which may be present in the full Lagrangian break R-invariance, and others do not. Scalar masses and self-interactions involving  $\phi^*\phi$  are invariant for any choice of the R-charge of the associated superfield. However R-invariance

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<sup>2</sup>See ref. [3] for a more detailed discussion, and the discussion of the vector superfields.

for the terms  $A\tilde{m}\tilde{t}\tilde{t}^c H_u$  and  $B\tilde{m}\mu H_u H_d$  are inevitably inconsistent with the conditions (1), since the  $\theta\theta$  which must be factored out in going from the superpotential to the Lagrangian written in terms of component fields carries  $R=2$ . Thus if either  $A$  or  $B$  is non-zero,  $R$ -invariance is broken explicitly<sup>3</sup>. However even if  $A = B = 0$ ,  $R$ -invariance is broken spontaneously when  $\langle H_u \rangle$  and  $\langle H_d \rangle$  are non-zero<sup>4</sup>, so that gaugino mass terms can be generated radiatively. It is interesting to consider several different possibilities, always taking tree-level gaugino masses to be zero:

1.  $A = B = 0$ . This corresponds to the possibility that  $R$ -invariance is only broken spontaneously, along with electroweak gauge invariance, by vevs of the Higgs fields.
2.  $A = 0$ . This corresponds to the absence of dimension-3 SUSY-breaking terms in the low energy Lagrangian, which arises naturally in hidden sector models without gauge singlets[4].
3. Non-zero  $A$  and  $B$ . We consider this for completeness, in case someday a SUSY-breaking mechanism is discovered which has this feature.

Years ago, the possibility of tree-level-vanishing gaugino masses in  $N=1$  supergravity theories was discussed in refs. [5](BGM) and [6](BM). These papers evaluated the leading radiative corrections to gaugino masses in a class of supersymmetric extensions of the standard model. Since then the world-view has changed considerably, because the top and Higgs are proving

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<sup>3</sup>If one chooses to define  $R$ -invariance to be the chiral symmetry associated with a massless gluino, with no reference to the transformation of  $\theta$  in the superfield, one would still arrive at the same conclusions, as a result of needing to give non-trivial transformations to quarks and squarks on account of their Yukawa couplings to gluinos and to Higgs and higgsinos on account of their Yukawa couplings to quarks and squarks.

<sup>4</sup>It is possible to find a solution to the conditions (1) such that either  $H_u$  or  $H_d$  has  $R = 0$ , but not both, so that if only one of the Higgs got a vev,  $R$ -invariance would not be broken spontaneously. Then it would be hard to understand ordinary fermion masses so we discard this as an option.

to be heavier than envisaged in those days, and because LEP constraints on new particles can be brought to bear. Furthermore the understanding of SUSY and ew symmetry breaking has advanced enough that much of the model-dependence of early work can be avoided. In this note we extend the BGM/BM analysis, eliminating recourse to a specific model of the symmetry breaking. In particular, we avoid their assumptions that  $A = 3$  and  $\mu = \tilde{m}$ , where  $\tilde{m}$  is the soft SUSY-breaking mass contribution common to all scalars. We generalize their results to arbitrary  $\tan\beta$  (the ratio of vev's of the two Higgs doublets which are responsible for electroweak symmetry breaking in supersymmetric models:  $\tan\beta \equiv \frac{v_u}{v_d}$ ). We also include radiative corrections to the chargino and neutralino mass matrices which have previously been neglected and which prove to be important in some regions of parameter space.

Two types of diagrams give the main radiative contributions:

1. Top-stop loops contribute to the gluino mass and to bino-w3ino ( $\tilde{b} \tilde{w}_3$ ) and  $\tilde{b} \tilde{b}$  entries in the neutralino mass matrix. The one loop contribution[5, 6] is proportional to the top mass times a function of the masses of the stop quark eigenstates  $m_{t1}$  and  $m_{t2}$ , which vanishes when they are degenerate. There can also be important 2-loop contributions coming from the top-stop loop with an additional Higgs exchange if  $A$  or  $B$  are non-zero.
2. One loop diagrams containing a W or Higgs and a wino, bino or higgsino contribute to the  $\tilde{w}_3 - \tilde{w}_3$ ,  $\tilde{b} - \tilde{b}$  and  $\tilde{w}_+ - \tilde{w}_+$  terms in the neutralino and chargino mass matrices. This contribution is proportional to  $\mu$ , the SUSY invariant coupling between the two Higgs superfields in the superpotential, times a function of tree level chargino and neutralino masses.

## 2 Experimental Constraints on Parameters

Since the masses of the charginos and squarks are constrained to be above about 45 GeV from their non-observation at LEP<sup>5</sup> the first step of our analysis is to express these masses in terms of the parameters  $\mu$ ,  $\tilde{m}$ ,  $A$ ,  $B$ , and  $\tan\beta$  of the theory, in order to determine which regions of parameter space are allowed in this scenario. In order to make our analysis independent of the details of the mechanism of ew symmetry breaking, we do not constrain the parameter space to guarantee the conditions for radiative electroweak symmetry breaking mechanism[1, 2].

### 2.1 Chargino Masses at Tree Level

We denote the chargino mass matrix by

$$\begin{pmatrix} 0 & X^T \\ X & 0 \end{pmatrix} \quad (2)$$

which acts on the spinor  $(\Psi^+ \ \Psi^-)$ , where  $\Psi^\pm$  are two component spinors:  $\Psi^+ = (\tilde{w}_+ \ \tilde{h}_u^+)$  and  $\Psi^- = (\tilde{w}_- \ \tilde{h}_d^-)$ . At tree level,  $X$  is the matrix

$$\begin{pmatrix} 0 & \sqrt{2} m_W \sin\beta \\ \sqrt{2} m_W \cos\beta & \mu \end{pmatrix}. \quad (3)$$

The parameter  $\mu$  does not violate supersymmetry. It enters the superpotential through the term  $\mu \hat{H}_u \hat{H}_d$ . We relate  $X$  to the diagonal matrix  $\mathcal{M} = UXU'$ , where  $U$  is the rotation

$$\begin{pmatrix} c & -s \\ s & c \end{pmatrix} \quad (4)$$

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<sup>5</sup>C.f. ref. [7]. We will also consider below the implications of lowering the  $m_{stop}$  bound to 15 GeV as discussed in ref. [8]. We avoid the use of CDF constraints on squark masses at this stage, since they depend on model dependent properties of other SUSY particles than the squarks[9].

and  $U'$  is obtained from  $U$  by  $c \rightarrow c', s \rightarrow -s'$ . In terms of the eigenvalues  $m_1, m_2$  of  $\mathcal{M}$ , the tree level chargino mass matrix, we have  $c \cdot c' = \mu \frac{m_2}{m_2^2 - m_1^2}$  and  $s \cdot s' = \mu \frac{m_1}{m_1^2 - m_2^2}$ . In order that the lighter eigenstate,  $m_2$ , is heavier than  $m_2^{lim}$ ,  $\mu$  must satisfy

$$\mu^2 < \frac{m_W^4 (\sin 2\beta)^2 - 2m_W^2 (m_2^{lim})^2 + (m_2^{lim})^4}{(m_2^{lim})^2}. \quad (5)$$

For  $m_2^{lim} = 45$  GeV and  $\beta = \frac{\pi}{4}$ , this gives  $\mu \lesssim 100$  GeV; the limit on  $\mu$  is lower for other choices of  $\beta$ . Fig. 1 shows the upper limit on  $\mu$  from the chargino mass limit, eqn (5), as a function of  $\beta$ .

## 2.2 Chargino Masses at One-Loop Level

Radiative corrections to the chargino mass matrix become significant if  $\mu$  is very large compared to  $m_W$ . Then the entries in  $X$  are modified by corrections which can be comparable to the off-diagonal elements in the tree level matrix. Taking  $\mu$  to be much larger than any other entry in  $X$ , the radiatively-corrected mass of the lighter chargino,  $m_2^{rc}$ , is essentially equal to the correction to the  $\tilde{w} \tilde{w}$  entry in (3). The main contribution arises at the one-loop level with charginos, neutralinos, gauge and Higgs bosons running in the internal lines. Its exact expression depends on the detailed mass spectrum of all these particles. However since we are interested in the large- $\mu$  limit, a major simplification occurs. The higgsinos,  $\tilde{h}_u$  and  $\tilde{h}_d$  combine together to form a Dirac SU(2) doublet of mass  $\mu$ , while the light eigenvectors of the chargino and neutralino mass matrices are mainly gauginos. A similar simplified pattern occurs also in the scalar Higgs sector. It is known[10] that, in the limit where one switches off the U(1) hypercharge coupling, the neutral Higgs scalar potential exhibits an SU(2)<sub>L</sub> × SU(2)<sub>R</sub> global invariance. The ew breaking breaks this global symmetry. However if  $\mu$  and/or  $\tilde{m}$  are  $\gg m_W$ , the corrections to the above global symmetry are small and we

can still classify the spin-0 mass eigenstates into two approximate SU(2) doublets. They are obtained from linear combinations of  $H_u$  and  $H_d$ , suitably weighted by  $\cos\beta$  and  $\sin\beta$  coefficients. One combination contains the massless SU(2)xU(1) would-be-Goldstone bosons and one light neutral Higgs (whose mass is  $\approx m_W$ ). Its couplings to the external gauginos are fixed by the Higgs mechanism. The orthogonal combination contains the heavy charged and neutral Higgs bosons, whose couplings are fixed by the orthogonality condition. If we denote by  $M$  the mass of these latter bosons,  $M$  will be of order of the larger of  $\mu$  and  $\tilde{m}$ .

Making use of the above mass spectrum, one obtains the following one-loop contribution to the  $\tilde{w}\tilde{w}$  entry in (3):

$$\delta_{\tilde{w}\tilde{w}} = \frac{\alpha_2}{2\pi} \frac{\mu m_1 m_2}{m_1^2 - m_2^2} \left( 3F(m_W, m_1, m_2) + \frac{m_2^2}{m_W^2} F(m_2, m_1, M) - \frac{m_1^2}{m_W^2} F(m_1, m_2, M) \right), \quad (6)$$

where

$$F(x, y, z) = \frac{z^2}{x^2 - z^2} \log\left[\frac{z^2}{x^2}\right] - \frac{y^2}{x^2 - y^2} \log\left[\frac{y^2}{x^2}\right]. \quad (7)$$

When  $\mu$  and  $M$  are much larger than  $m_2 \lesssim m_W$ , this gives approximately

$$m_2^{rc} = \frac{\alpha_2}{\pi} \cos\beta \sin\beta \frac{\mu M^2}{(\mu^2 - M^2)} \log\left[\frac{\mu^2}{M^2}\right]. \quad (8)$$

This expression shows that the present experimental bound on the lighter chargino can be accommodated for sufficiently large  $\mu$  and  $M$ . As noted above, when  $\tilde{m} \gg \mu$ , we expect  $M \sim \tilde{m}$ , while if  $\tilde{m} \ll \mu$  we expect  $M \sim \mu$ . Simply using eqn (8) with  $M = \max[\tilde{m}, \mu]$  gives the boundaries of the allowed regions of  $\mu$  (on the horizontal axis) and  $\tilde{m}$  (on the vertical axis) shown in Fig. 2, for  $\beta = \frac{\pi}{4}$ . Once  $\mu \gtrsim \frac{\pi m_2^{lim}}{\alpha_2 \cos\beta \sin\beta}$ , the condition (8) is satisfied for any  $\tilde{m}$ , accounting for the vertical segments. For the present value of  $m_2^{lim} = 45$  GeV we have the dashed curve, while the dot-dashed curve gives the boundary of the allowed region for  $m_2^{lim} = 80$  GeV. The allowed region is above and to the right of these curves, but remember that the sharp corners and

vertical lines are artifacts of the simplistic relation  $M = \max[\tilde{m}, \mu]$ . Given a particular model, one can find the smooth curve which this approximates. For other values of  $\beta$ , the boundary shown in the figure should be multiplied by  $[2\cos\beta\sin\beta]^{-1}$ . For  $\mu \lesssim 2$  TeV one sees that the  $M$  required becomes very large: for  $\mu = 2$  (1.5, 1) TeV,  $M$  must be larger than 16 (27, 94) TeV, respectively. Thus in this tree-level-massless gaugino scenario, unless one is willing to consider very large values of  $M$  there is effectively a gap in allowed  $\mu$ 's between  $\sim 100$  GeV and a few TeV.<sup>6</sup> If no chargino is found at LEP II and  $m_2^{lim}$  is increased to above  $m_W$ , then the entire low- $\mu$  region will be removed and the allowed range of  $\mu$  and  $M$  is just the area above the dot-dashed curve in Fig. 2.

## 2.3 Stop Mass

The top squark  $mass^2$  matrix for the effective low-energy theory is approximately<sup>7</sup>:

$$\begin{pmatrix} m_t^2 + \tilde{m}^2 + m_Z^2 \cos(2\beta) \left(\frac{1}{2} - \frac{2}{3} \sin^2 \theta_w\right) & A_{\text{eff}} m_t \tilde{m} + \mu m_t \cot \beta \\ A_{\text{eff}} m_t \tilde{m} + \mu m_t \cot \beta & m_t^2 + \tilde{m}^2 + m_Z^2 \cos(2\beta) \left(\frac{2}{3} \sin^2 \theta_w\right) \end{pmatrix}, \quad (9)$$

where

$$A_{\text{eff}} \equiv A + \delta A + B \frac{k\mu^2}{M^2} \quad (10)$$

and  $\delta A$  is the radiative correction to  $A$  from the gluino-top loop, in which the gluino mass insertion is a one-loop diagram. When  $A = B = 0$  at tree level this correction can be relevant for some regions of parameters. From formulae

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<sup>6</sup>If the tree-level mass of charginos is of O(50-90) GeV, as studied in the recent preprint [11], then the radiative corrections considered here can be large enough to be experimentally significant for smaller values of  $\mu$ . GF thanks M. Strassler for a discussion of this work.

<sup>7</sup>The  $mass^2$  matrix (9) takes the low-energy soft-susy-breaking contributions to the  $t\tilde{t}$  and  $\tilde{t}^c \tilde{t}^c$  entries equal to a single parameter  $\tilde{m}^2$ . Modification of this simplest assumption, e.g., due to RG running, is discussed below.



given in ref. [1], its magnitude is  $\delta A = \frac{4\alpha_s}{\pi} \frac{m_{\tilde{g}}}{\tilde{m}} \log(\frac{M_{initial}}{m_W})$ . If  $B \neq 0$ , a  $\tilde{t}\tilde{t}^c$  mixing can also arise from the vertices  $h_t \mu \tilde{t}\tilde{t}^c H_d$  and  $B \tilde{m} \mu H_u H_d$  connected by an  $H_d$  propagator. Evaluating this propagator at zero momentum and taking the vev of  $H_u$  produces a contribution to  $A_{\text{eff}}$  proportional to  $B$ . Its sign and precise magnitude depend on the details of the Higgs mass spectrum, so we parameterize it in terms of  $k$ , a constant which is presumably of order one, and a generic scalar Higgs mass,  $M$ .

When dimension-3 SUSY-breaking operators are absent from the low energy theory,  $A = 0$ ; more commonly it has been taken to be of order 1, e.g., 3 in BM[6].  $A_{\text{eff}}$  cannot be made too large or the scalar quarks or leptons will get a vev and color SU(3) or electromagnetism will be broken. Typically this leads to an upper bound on the modulus of  $A$  close to 3[12, 1, 13, 14]. We will see below that consistency with the experimental lower bound on the lighter stop mass generally requires  $A_{\text{eff}}$  to be even smaller than this.

The diagonal terms in the stop  $mass^2$  matrix determine the average squark  $mass^2$ . Splitting between the physical stop mass eigenstates is mainly controlled by the off-diagonal terms, as long as the diagonal terms are not too different. Thus an experimental lower limit on the stop mass,  $m_{stop}^{lim}$ , implies an upper limit on  $A_{\text{eff}} \tilde{m} + \mu \cot \beta$  for a given average stop mass-squared. Dropping the small  $m_Z^2$  corrections in (9) to make the point clear, this is

$$A_{\text{eff}} \tilde{m} + \mu \cot \beta \lesssim \frac{\tilde{m}^2 + m_t^2 - (m_{stop}^{lim})^2}{m_t}. \quad (11)$$

From this expression one sees that the limit on  $\mu$  from the stop mass constraint is essentially independent of  $m_{stop}^{lim}$  and  $A_{\text{eff}}$  when  $\tilde{m}$  is large. The upper limit on  $\mu$  for a given  $\tilde{m}$  is shown in the large  $\mu$  region as the solid line Fig. 2; the allowed region of  $\tilde{m}$  for a given  $\mu$  is above the line. One sees that in the large  $\mu$  region if the chargino constraint is satisfied, the squark constraint will usually be also.

In the small  $\mu$  region when  $\tan \beta \geq 1$ , the constraint from the stop mass limit is less stringent than from the chargino limit, except for very small  $\tilde{m}$ .

Since the CDF limits on squark masses must be reexamined when the gluino becomes as light as we will be considering, we use  $m_{stop}^{lim} = 45$  GeV to be conservative. However even if the strongest CDF limit of  $m_{stop}^{lim} = 126$  GeV were applicable, we found that it would make an insignificant difference in these limits except for  $A_{eff} \neq 0$  and small  $\tilde{m}$ . We have checked that modifying the  $\tilde{m}^2$  terms in the diagonal elements of (9) as would arise from different renormalization group running of the  $\tilde{t}$  and  $\tilde{t}^c$  masses in the RG-induced ew symmetry breaking scenario, does not significantly affect these conclusions, again because the chargino mass provides the more stringent constraints on parameters.

Thus for most of the interesting parameter space in the small as well as large  $\mu$  region, consistency with the LEP chargino and squark mass limits is guaranteed simply by satisfying eqn (5) from the chargino limit, independent of the stop mass limit,  $\tilde{m}$ , and  $A_{eff}$  (as long as it is not too large). Note however that for larger  $A_{eff}$  the stop mass limit becomes dominant and in fact requires that  $A_{eff}$  be less than some maximum value for given  $\tilde{m}$  and stop mass limit. Fig. 3 shows this, for  $\tilde{m} = 100$  (solid), 250 (dashed), and 400 GeV (dot-dashed). The upper plot uses the stop  $mass^2$  matrix (9), while for the lower plot the  $\tilde{m}^2$  in the (1,1) and (2,2) element of (9) has been modified to  $\frac{2}{3}\tilde{m}^2$  and  $\frac{1}{3}\tilde{m}^2$ , respectively. This simulates (see Table 1 of ref. [1]) the case that these terms are equal at the susy-breaking scale but experience RG running which also causes the ew gauge symmetry to break.

To summarize, requiring the lightest SUSY charged particles to be heavier than the experimental lower bounds leads to two distinct allowed regions for  $\mu$  and associated regions for  $\tilde{m}$  – namely  $\mu \lesssim 100$  GeV, or  $\mu \gtrsim$  several TeV. Now let us find the gluino and lightest neutralino ( $\chi_1^0$ ) masses for the allowed parameter regions.

### 3 The Gluino

The top-stop loop produces the only important 1-loop correction to the gluino mass[5, 6]:

$$\delta_g^{(1)} = \frac{\alpha_s m_t}{4\pi} \sin(2\theta_t) F(m_t, m_{t1}, m_{t2}), \quad (12)$$

where the function  $F$  is the same as in eqn (7) and  $\theta_t$  is the rotation which diagonalizes the stop mass matrix<sup>8</sup>. Using (9) for the stop  $mass^2$  matrix,

$$\sin^2(2\theta_t) = \frac{[(A_{\text{eff}}\tilde{m} + \mu\cot\beta)m_t]^2}{[(A_{\text{eff}}\tilde{m} + \mu\cot\beta)m_t]^2 + \frac{1}{4}[m_Z^2\cos(2\beta)(\frac{1}{2} - \frac{4}{3}\sin^2\theta_W)]^2}. \quad (13)$$

In this case,  $\sin(2\theta_t) \approx 1$  for most of parameter space. We will consider below the case that (9) is modified such that the soft-susy breaking contributions to the diagonal elements of the stop  $mass^2$  matrix are not equal. Note that  $F(x, y, z)$  is odd under  $y \leftrightarrow z$  so that  $\delta_g^{(1)}$  can be seen to vanish linearly with the fractional splitting between the stop mass eigenstates. Having  $A_{\text{eff}}$  non-zero or having a large value of  $\mu\cot\beta$  contributes to a larger gluino mass because each of these increases the mass splitting between stop mass eigenstates (see eqn (9)).

The top-stop contribution to gaugino masses can have a 2-loop divergent piece coming from Higgs exchange between top and stop, if the dimension-3 SUSY-breaking scalar trilinear coupling  $A\tilde{m}\tilde{t}\tilde{t}^c H_2$  is non-vanishing. All divergent 2-loop diagrams have been calculated recently in refs. [15, 16] and we use their result here.<sup>9</sup> Denoting by  $M_{\text{initial}}$  the renormalization scale at which the counterterm exactly cancels the contribution of this divergent graph, so that gauginos are massless, the RG contribution to the low energy

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<sup>8</sup>We thank D. Pierce for pointing out that we had omitted writing this factor in the original version of the manuscript; it has been included in the numerical analysis.

<sup>9</sup>The newer results differ by an overall factor of 3 from that given in [6]. Note that [6] corrects a factor-of-2 error in the one-loop contribution given in [5].

gluino mass is:

$$\delta_g^{(2)} = \frac{\alpha_3 \alpha_2 (1 + \cot^2 \beta)}{4\pi^2} A \tilde{m} \left( \frac{m_t}{m_W} \right)^2 \log \left( \frac{M_{initial}^2}{m_W^2} \right). \quad (14)$$

In addition, if  $B$  is non-zero, there is a finite 2-loop contribution to the gluino mass which can be important for some portions of parameter space. It corresponds to the same diagram as the one just considered, but with the pointlike vertex  $A \tilde{m} \tilde{t} \tilde{t}^c H_u$  replaced by the vertices  $h_t \mu \tilde{t} \tilde{t}^c H_d$  and  $B \tilde{m} \mu H_u H_d$  connected by an  $H_d$  propagator. For  $\mu$  and  $\tilde{m} \gg m_W$  and  $m_t$ , this correction is approximately

$$\delta_g^{(3)} \approx \frac{\alpha_3 \alpha_{em}}{4\pi^2 \sin^2 \theta_W} \frac{B \mu^2 \tilde{m}}{\bar{M}^2} \left( \frac{m_t}{m_W} \right)^2, \quad (15)$$

where  $\bar{M}$  denotes the highest mass in the loop. Taking  $\bar{M} = \max[\mu, \tilde{m}]$ , this contribution is maximized for  $\mu = \tilde{m}$ , for which it is  $4 \times 10^{-4} \mu B$ . Thus it is only relevant if  $A = 0$  and  $\mu \sim \tilde{m}$ , with  $\mu B$  of order several TeV or larger.

The first column of Fig. 4 shows the one-loop contribution to the gluino mass,  $\delta_g^{(1)}$ , for  $A_{\text{eff}} = 0$  and  $A_{\text{eff}} = 1$ , as a function of  $\mu$  in the low  $\mu$  region, for  $\beta = \frac{\pi}{4}$  and several choices for  $\tilde{m}$ . Also in the low  $\mu$  region, the first column of Fig. 5 shows the gluino mass as a function of  $\beta$  at the maximum value of  $\mu$  which is consistent with whichever is the stronger of the stop mass or chargino mass constraints (in fact, almost always the latter), for  $A_{\text{eff}} = 0$  and 1. These results are computed with the stop  $mass^2$  matrix (9). If there are significant differences in the low-energy soft-susy-breaking diagonal terms in the stop  $mass^2$  matrix, the gluino mass predictions are modified somewhat. To illustrate the possible extent of this effect, consider the scenario of radiative ew symmetry breaking. In that case, the  $\tilde{m}^2$  in the 1, 1 and 2, 2 elements of (9) is multiplied by  $\sim 2/3$  and  $\sim 1/3$  respectively, taking the values of the corrections chosen in the previous section as an example. For  $\tilde{m} \gtrsim m_t$  the change is quantitatively although not qualitatively important. We give the 1-loop contributions to the gluino mass predictions for this case in Fig. 6.

For  $\tilde{m} \gtrsim m_t$ , the 1-loop contribution to the gluino mass decreases as  $\tilde{m}$  is increased with  $A$  and  $\mu$  held fixed. This is because increasing  $\tilde{m}$  decreases the fractional splitting between the stop mass eigenstates,  $\sim \frac{(A_{\text{eff}}\tilde{m} + \mu \cot\beta)m_t}{(\tilde{m}^2 + m_t^2)}$ . Thus for  $A = B = 0$  the gluino mass is negligible in the large  $\mu$ ,  $\tilde{m}$  region, unless  $\mu m_t \sim \tilde{m}^2$ . The maximum value of the gluino mass in this latter case occurs when the lighter stop is as light as is allowed experimentally while the heavier stop is very massive, thus maximizing the fractional splitting between eigenstates. Figure 7 shows the maximum gluino mass under these circumstances, with  $m_{\text{stop}}$  greater than 45 (dashed) and 126 (dot-dashed) GeV. The relationship required to implement this,  $\mu \approx \tilde{m}^2/m_t \gg \tilde{m}$ , is unconventional.

If  $\tilde{m}$  is large and  $A \neq 0$  the divergent 2-loop contribution can be important. Fig. 8 shows  $\delta_g^{(2)}$  for  $A = 1$ ,  $B = 0$  and  $\beta = \frac{\pi}{4}$ . The solid curve corresponds to taking  $M_{\text{initial}} \rightarrow \tilde{m}$ , giving an estimate of the minimal importance of this correction, while the dashed and dot-dashed curves show the result for  $M_{\text{initial}} = 10^{11}$  GeV and  $M_{\text{initial}} = M_{\text{pl}}$ . Evidently, for large  $M_{\text{initial}}$  this is a large effect. It would be interesting to determine the value of this two-loop contribution imposing as well the constraints of the radiative electroweak breaking scenario.

To summarize, for the 3 cases we are treating,

1.  $A = B = 0$ : In the low  $\mu$  region the gluino mass decreases with increasing  $\tilde{m}$  from  $\lesssim 700$  MeV for  $\tilde{m} = 100$  GeV to  $\lesssim 200$  MeV for  $\tilde{m} = 400$  GeV (Fig. 5, upper left, and Fig. 6, upper right). In the large  $\mu$  region the gluino mass is negligible unless  $\mu \sim \tilde{m}^2/m_t$ , in which case the maximum gluino mass is  $\sim 6$  GeV for  $\mu \lesssim 20$  TeV (Fig. 7).
2.  $A = 0$ ,  $B \neq 0$ : When  $\mu \ll \tilde{m}$ , this case is equivalent to the previous case with  $A = B = 0$ . For  $\mu \sim \tilde{m}$  in the low  $\mu$  region, the  $kB\frac{\mu^2}{M^2}$  contribution to  $A_{\text{eff}}$  can produce  $A_{\text{eff}} \sim 1$  so that gluino masses can be of order a few GeV (see Fig. 5, lower left plot, and Fig. 6, lower right

plot). For the large  $\mu \sim \tilde{m}$  region the two loop diagram proportional to  $B$  makes a contribution (eqn 15)  $\sim 4 \times 10^{-4} \mu B$ .

3.  $A \neq 0$ : In the low  $\mu$  region this gives gluino masses of order a few GeV as discussed in the item above. However in the large  $\mu$  region the gluino mass can be very large due to the 2-loop divergent diagram: e.g., for  $A = 1$  and  $M_{initial} \gtrsim 10^{11}$  GeV, the gluino mass is consistent with the present CDF missing energy bound[9] as long as  $\mu \gtrsim 8$  TeV (see Fig. 8).

## 4 The Lightest Neutralino

The tree level neutralino mass matrix, in the basis  $(\tilde{b}, \tilde{w}_3, \tilde{h}_1, \tilde{h}_2)$ , is:

$$\begin{pmatrix} 0 & 0 & -M_Z \cos \beta \sin \theta_w & M_Z \sin \beta \sin \theta_w \\ 0 & 0 & M_Z \cos \beta \cos \theta_w & -M_Z \sin \beta \cos \theta_w \\ -M_Z \cos \beta \sin \theta_w & M_Z \cos \beta \cos \theta_w & 0 & -\mu \\ M_Z \sin \beta \sin \theta_w & -M_Z \sin \beta \cos \theta_w & -\mu & 0 \end{pmatrix}. \quad (16)$$

Radiative corrections remove the zeros in this matrix. Let us first consider the radiative contributions to the neutral gaugino 2x2 sub-matrix in the upper left-hand corner.

The  $\tilde{b} - \tilde{w}_3$  off-diagonal entries receive one- and two-loop contributions entirely analogous to those that we computed for the gluino mass:

$$\delta_{\tilde{b}\tilde{w}} = \frac{\sqrt{\alpha_1 \alpha_2}}{\alpha_3} m_{\tilde{g}}, \quad (17)$$

where  $m_{\tilde{g}} = \delta^{(1)} + \delta^{(2)} + \delta^{(3)}$ , given in eqns (12),(14) and (15). As for the diagonal entries, the contribution to  $\tilde{w}_3 - \tilde{w}_3$  is readily related to  $\delta_{\tilde{w}\tilde{w}}$  in the chargino sector, eqn (6), in the approximation of large  $\mu$  that we discussed there. Finally, the  $\tilde{b}\tilde{b}$  entry receives two types of radiative contributions. The

first comes from one- and two-loop corrections with top and stop running in the loops, yielding a contribution proportional to  $m_{\tilde{g}}$  analogous to the expression in eqn (17). The other type of correction is from higgsino-higgs loops<sup>10</sup>. It is the same as for  $\tilde{w}\tilde{w}$ , replacing  $\alpha_2$  by  $\alpha_1$ . All together we obtain:

$$\delta_{\tilde{b}\tilde{b}} = \frac{2\alpha_1}{3\alpha_3}m_{\tilde{g}} + \frac{\alpha_1}{\alpha_2}\delta_{\tilde{w}\tilde{w}}. \quad (18)$$

We do not compute the radiative corrections to the higgsino submatrix in detail, since they depend on the model of ew symmetry breaking. They would not be present if there were a Peccei-Quinn symmetry, so that they must be proportional to  $\mu$  and/or the vev's of  $H_u$  and  $H_d$ . We find that these radiative corrections cannot be larger than the  $O(\mu)$  entries which are present in eqn (16). We checked that the masses of the lightest two neutralinos change only slightly when such terms are included.

We find the eigenvalues and eigenvectors of the radiatively-corrected neutralino mass matrix numerically, for a variety of values of parameters. The mass of the lightest neutralino,  $\chi_1^0$ , in the small  $\mu$  region is shown in the second column of Figs. 4 and 5 as a function of  $\mu$  and  $\beta$ , respectively. In the small  $\mu$  region,  $m(\chi_1^0)$  is rather insensitive to  $\tilde{m}$  and  $A_{\text{eff}}$ , but is sensitive to  $\mu$ . Typical values for  $m(\chi_1^0)$  in the small- $\mu$  case are  $O(\frac{1}{10} - 1)$  GeV. Since the top-stop loop is responsible for only a fraction of the neutralino mass,  $m(\chi_1^0)$  is insensitive to possible differences in the susy-breaking diagonal stop  $mass^2$  and the analog of Fig. 6 is not needed. The dependence of  $m(\chi_1^0)$  on  $\mu$  and  $\tilde{m}$  in the large  $\mu$  region is shown in Fig. 9. There one sees that  $m(\chi_1^0) \gtrsim 10$  GeV for the large- $\mu$  case. The sensitivity to  $A_{\text{eff}}$  is too small to be seen in this scale figure. The second-lightest neutralino,  $\chi_2^0$ , typically has a mass of about 50 GeV.

The gluino mass is more strongly dependent on  $A_{\text{eff}}$  than are the neutralino masses just because the gluino gets its mass entirely from quark

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<sup>10</sup>In Feynman 'tHooft gauge, where diagrams with gauge bosons in the loop vanish.

squark loops which are sensitive to  $A_{\text{eff}}$  and proportional to the gauge coupling constant appropriate to the gaugino in question. On the other hand, the bino-wino submatrix of the full neutralino mass matrix, whose eigenvalues are dominantly important in determining the lightest neutralino mass, is approximately

$$\begin{pmatrix} \frac{1}{15}m_{\tilde{g}} + \frac{1}{4}m_2^{rc} & \frac{1}{5}m_{\tilde{g}} \\ \frac{1}{5}m_{\tilde{g}} & m_2^{rc} \end{pmatrix}, \quad (19)$$

when the known gauge couplings are inserted into eqns (17) and (18). Evidently, the eigenvalues of (19) are insensitive to the top-stop loops unless the radiatively generated gluino mass is  $\gtrsim 4m_2^{rc}$ . Since we are only considering parameter ranges such that  $m_2^{rc} > 45$  GeV, the lightest neutralino mass is generically insensitive to  $A_{\text{eff}}$  unless  $m_{\tilde{g}} \gtrsim 200$  GeV. One can also see from (19) how restricting the parameter space further by improving the chargino mass limits would in general simply scale up the predictions for the masses of the lightest neutralinos in proportion to the chargino mass limit.

For parameters such that  $m(\chi_1^0) \lesssim 2$  GeV the composition of  $\chi_1^0$  and  $\chi_2^0$  are insensitive to parameters and

$$|\chi_1^0\rangle \approx 0.88|\tilde{b}\rangle + 0.47|\tilde{w}_3\rangle. \quad (20)$$

This is very close to the  $U(1) \times SU(2)$  composition of the photon, so in the small  $\mu$  region, the lightest neutralino is essentially a photino. The  $\tilde{b}$  component becomes more dominant with increasing  $m(\chi_1^0)$ , reaching about 0.99 for the large  $\mu$  scenario. In all cases, however, the higgsino components have amplitudes less than 1% for both  $\chi_1^0$  and  $\chi_2^0$ . This explains the insensitivity of the masses of the two lightest neutralinos to the model-dependent radiative corrections to the higgsino mass submatrix noted in the previous paragraph. Since the  $Z^0$  only decays to neutralinos through their higgsino components, the relative probability of a  $Z^0$  decaying to a pair of neutralinos, compared to decaying to a given neutrino-antineutrino pair, is  $\lesssim 10^{-8}$ . Thus the impressive experimental constraint from LEP on the number of extra neutrinos



is insufficient to limit the existence of these neutralinos.

## 5 Phenomenology and Cosmology

Now we briefly turn to the phenomenological viability of the scenario we have investigated. While our analysis above was general enough to include arbitrary  $A$ , it is particularly interesting to consider  $A = 0$ . This is because in hidden sector dynamical SUSY breaking without gauge singlets, all dimension-3 SUSY-breaking operators in the low energy theory, including a gaugino mass term and the trilinear squark-squark-Higgs coupling whose coefficient is defined to be  $A\tilde{m}$ , are suppressed by a factor  $\frac{\tilde{m}}{M_{pl}}$  and thus are expected to be very small<sup>11</sup>. As long as  $A_{\text{eff}}$  is small, the lightest neutralino is generically heavier than the gluino. For instance for  $\tan\beta = 1$  and  $\mu = 100$  GeV, the lightest neutralino mass falls in the range  $0.5 - 0.8$  GeV, while the gluino mass is found to be less than 0.3 GeV (see Fig. 4). In the large  $\mu$  region the lightest neutralino mass is greater than 10 GeV. Throughout the large  $\mu$  region the upper limit on the gluino mass consistent with the experimental lower limit on the stop and chargino masses is less than the lightest neutralino mass<sup>12</sup>.

The phenomenology of hadrons containing light gluinos is discussed in ref. [17] and references cited therein. Some essential conclusions are the following:

1. The theoretical lower limit on the gluino mass coming from requiring that the  $\eta'$  be a pseudogoldstone boson is  $m_{\tilde{g}} \sim 10 \frac{\langle q\bar{q} \rangle}{\langle \lambda\lambda \rangle} m_s$  [17]. The gluino condensate is very uncertain but is expected to be larger than the quark condensate. Conceivably the ratio is large enough to cancel

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<sup>11</sup>See ref. [4] for a more detailed discussion of the argument.

<sup>12</sup>Unless  $\mu$  and  $\tilde{m}$  are *very* large so that  $B$  terms can dominate.

the factor of 10, leading to a lower bound on the gluino mass of order one to several hundred MeV. This is just the range found above in the low  $\mu$  region for  $A_{\text{eff}} = 0$ , so that improvements in the determination of the  $\eta'$  mass as a function of the mass of a light gluino will allow part of the parameter space to be excluded. For  $A = B = 0$ , one can already exclude  $\tilde{m} \gtrsim 300$  GeV when  $\mu \lesssim 100$  GeV.

2. The non-observation[18] of any peak in the photon spectrum in radiative  $\Upsilon$  decay excludes gluinos in the mass range  $\sim 1.5 - 3.5$  GeV, for any lifetime. This excludes small regions of parameter space in the large  $\mu$  region.
3. Light gluinos would be mainly found in the flavor-singlet hadron  $R^0$ , a gluon-gluino bound state, or the flavor-singlet baryon  $S^0$  composed of  $uds\tilde{g}$ . The mass of the  $R^0$  can be estimated[17] from the lattice calculation of the mass of the  $0^{++}$  glueball to be  $1440 \pm 375$  MeV for a massless gluino.  $R^0$ 's with mass  $\lesssim 2.2$  GeV are experimentally allowed, except for lifetimes in the  $\sim 2 \times 10^{-6} - 10^{-8}$  sec range[19] or shorter than  $\sim 5 \times 10^{-11}$  sec, where beam dump experiments are useful[17] if  $\tilde{m}$  is not too large. In the small  $\mu$  scenario the lightest neutralino is typically heavier than the gluino, and the  $R^0$  decay rate is suppressed compared to the conventional phenomenological treatment in which the lightest neutralino is assumed to be essentially massless. Suitable methods to estimate the  $R^0$  lifetime must be developed to see if the present experimental limits constrain this scenario.
4. Long lived or absolutely stable  $R^0$  and  $S^0$  are not obviously excluded. They would not bind to nuclei, so would not be found in searches for exotic isotopes[17]. In fact, they could help provide the dark matter of the universe and might account for anomalous production of muon events by cosmic rays coming from Cygnus X-3[17]. Stable or very

long-lived  $R^0$  and  $S^0$ 's are practically assured in the large  $\mu$  region if  $A = 0$  because then they are lighter than the lightest neutralino.

Since short lived gluinos ( $\tau < 2 \times 10^{-11} \frac{m_{\tilde{g}}}{1\text{GeV}}$  sec) with masses between  $\sim 4 - 126$  GeV are excluded by missing energy searches (see ref. [17] for discussion and references), we can restrict the large  $\mu$  parameter space for  $A \sim 1$  by requiring[17]

$$\tau_{\tilde{g}} = \frac{128\pi \cos^2 \theta_W}{\alpha_s \alpha_{em}} f\left[\frac{m_\chi}{m_{\tilde{g}}}\right] \frac{M_{sq}^4}{m_{\tilde{g}}^5} > 2 \times 10^{-11} m_{\tilde{g}} \frac{\text{sec}}{\text{GeV}}, \quad (21)$$

where  $f[y]$  is the phase space suppression when the lightest neutralino mass is a non-negligible fraction of the gluino mass;  $f[0] = 1$ . We replace  $f \rightarrow 1$  to get a rough estimate. Then inequality (21) requires that either  $m_{\tilde{g}} \gtrsim 126$  GeV or

$$m_{\tilde{g}} \lesssim 28 \text{ GeV} \left( \frac{\tilde{m}}{10\text{TeV}} \right)^{2/3} f^{1/6}. \quad (22)$$

The latter condition requires that

$$A(1 + \cot^2 \beta) \log \left( \frac{M_{initial}^2}{m_W^2} \right) \lesssim 6.6 \left( \frac{\tilde{m}}{10\text{TeV}} \right)^{-1/3} f^{1/6} \quad (23)$$

be satisfied<sup>13</sup>.

Fully studying the constraints on this scenario coming from requiring relic particles not to overclose the universe is beyond the scope of this paper. In the usual scenario with  $A \sim 1$ , and tree level gaugino masses taken to be proportional to the squark masses, these considerations are used to rule out the existence of stable neutralinos having mass less than a few GeV<sup>14</sup>. However the contribution of a relic to the present mass density of the universe

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<sup>13</sup>The correction to these limits coming from retaining the phase space factor  $f$  is small unless  $\tilde{m} < 10$  TeV or  $\mu$  and  $\tilde{m}$  are very large. For instance, with  $m_{\tilde{g}} \sim 28$  GeV and  $m_\chi \sim 10$  GeV,  $f[\frac{10}{28}] = 0.56$  reducing the limit in comparison to the rough estimate by 10%.

<sup>14</sup>See, e.g., ref. [20, 21].

$\sim \frac{1}{\langle \sigma_{annih} v \rangle}$ , and is therefore more weakly dependent on the relic mass than on the squark mass because  $\sigma_{annih} \sim M_{sq}^{-4}$ . Furthermore, when the gluino is light the availability of the reaction  $\chi\tilde{g} \rightarrow q\bar{q}$  enhances the annihilation of the neutralino, because the cross section is larger by a factor  $\sim \frac{\alpha_s}{\alpha_{em}}$  and because, unlike  $\chi\chi$  annihilation, it can go via the s-wave so the cross section is non-vanishing in the non-relativistic limit[22]. Thus annihilation of neutralinos is more efficient in this scenario even for the same squark mass and, more importantly, limiting the squark mass puts different constraints on the gluino mass than in the usual scenario. For the large  $\mu$  region this can be analysed without difficulty. However when the gluino and photino are in the  $\lesssim 1$  GeV range, the freeze-out temperature is of the same order of magnitude as the QCD confinement phase transition temperature, so that the discussion of this scenario is considerably more complicated than in the usual case and detailed analysis is required to make quantitative statements.

Since a chargino has not been seen at LEP, we inferred in Section 2 that  $\mu$ , the supersymmetric coupling between the two Higgs doublets, is either less than  $\sim 100$  GeV or greater than several TeV in this scenario. If a chargino is not discovered at LEP II, the low  $\mu$  region would also be excluded. If it were possible to exclude the large  $\mu$  region on other grounds, this would mean that the present scenario could be definitively excluded at LEP II. We have not made a comprehensive study of other constraints on  $\mu$ , but note that for a given model of ew symmetry breaking only certain regions of  $\mu$  will be allowed. For instance, the radiative breaking scenario as discussed in refs. [1, 2] does not work in the large  $\mu$  region when  $A = 0$ .

## 6 Summary

We have investigated radiative corrections to gaugino masses, revealing a number of interesting new possibilities for the gaugino sector of a supersym-

metrized standard model. Constraining the parameters of the model so that the lightest supersymmetric charged particles are consistent with experimental bounds, we find that if R-invariance is only broken spontaneously or if the dimension-3 SUSY-breaking parameters which explicitly violate R-invariance are absent, the lightest neutralino is typically heavier than the gluino. In the low  $\mu$  region, the masses of the gluino and lightest neutralino are less than  $\sim 2$  GeV, even when  $A$ , the dimension-3 squark-squark-Higgs coupling, is non-zero. In the large  $\mu$  region the lightest neutralino is heavier than  $\sim 10$  GeV and is more massive than the gluino unless  $A \neq 0$ . Thus the lightest gluino-containing hadron naturally tends to be long-lived or even stable, and can be consistent with laboratory searches[17]. While this scenario is very unusual from the phenomenological and cosmological points of view, it may be consistent with observations. Further work is needed to constrain the parameters of the model from considerations other than just charged particle masses, and to explore the experimental and cosmological implications of this scenario in greater detail<sup>15</sup>. A more complete discussion of these issues is left to the future<sup>16</sup>.

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<sup>15</sup>A discussion of various astrophysical consequences of a stable gluino can be found in a recent preprint by Plaga[23].

<sup>16</sup>*Note Added:* We wish to thank D. Pierce for calling our attention to his paper with A. Papadopoulos[24] which deals with some of the issues we discuss here. They also determine the radiative corrections to chargino and neutralino masses. In principle the case we treat should be obtainable as a special case of their formulae, however our expressions are considerably more compact and transparent than theirs and we have not attempted to make a comparison. Their work is complementary to ours, in that it focuses on the possibility of extracting information on the GUT-scale mass relations from observed sparticle masses, assuming general tree-level gaugino masses. By virtue of their interest in generality, they did not explore in detail the scenario which we find most interesting, namely the possible absence of dimension-3 susy-breaking terms. The consequences of this form of susy-breaking for the phenomenologically crucial issue of the relative masses of gluino and lightest neutralino is the main new feature of the present work. Specializing to the portions of their discussion relevant to a massless gaugino scenario, one finds that the regions of parameter space considered acceptable in ref. [24] differ from ours in important ways. For instance we find that for tree-level massless gauginos, the lower limit on the chargino mass severely restricts the  $\mu - \beta$  space, so the large values of  $\tan\beta$  which they consider (see

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e.g., their Figs. 2 and 4) are actually excluded in this scenario, at least as long as we are in a region of parameter space where our global  $SU(2)$  ansatz holds. Another important difference is that in their discussion of the neutralino sector, they seem to consider a tree-level neutralino mass to be a necessity as a result of using LEP and CDF limits which are, however, not applicable when one considers the very specific phenomenology which follows from the absence of dimension-3 SUSY-breaking as described here and in ref. [17].

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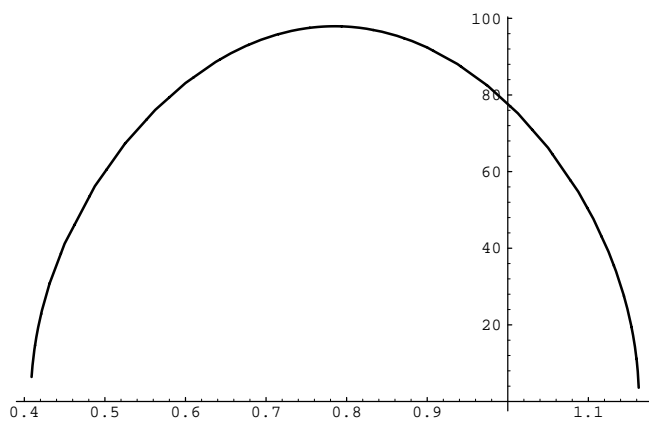


Figure 1: Maximum allowed value of  $\mu$  in GeV as a function of  $\beta$  from the chargino mass limit.

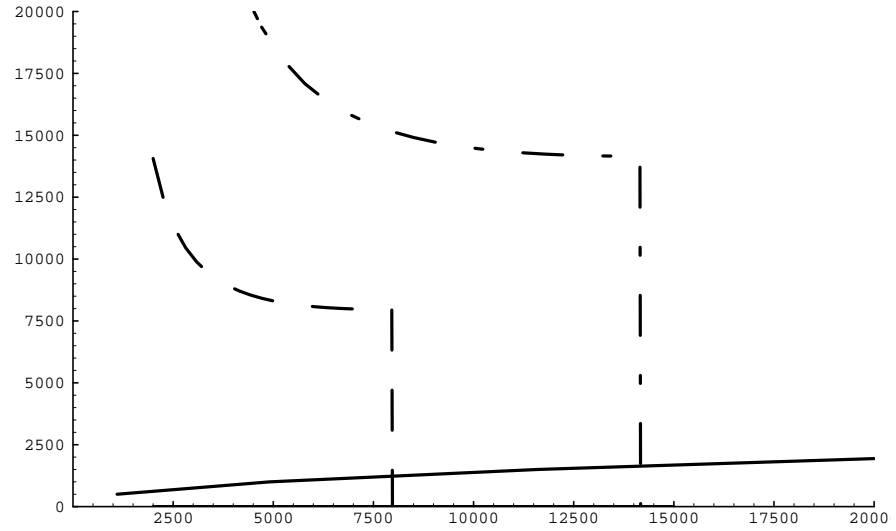


Figure 2: Lower limits on  $M$  as a function of  $\mu$  for consistency with a chargino mass limit of 45 (dashed) and 80 (dot-dashed) GeV. Solid curve is the boundary of the region allowed by the squark mass limit, with the allowed region being above the curve. All masses are in GeV in this and other figures.

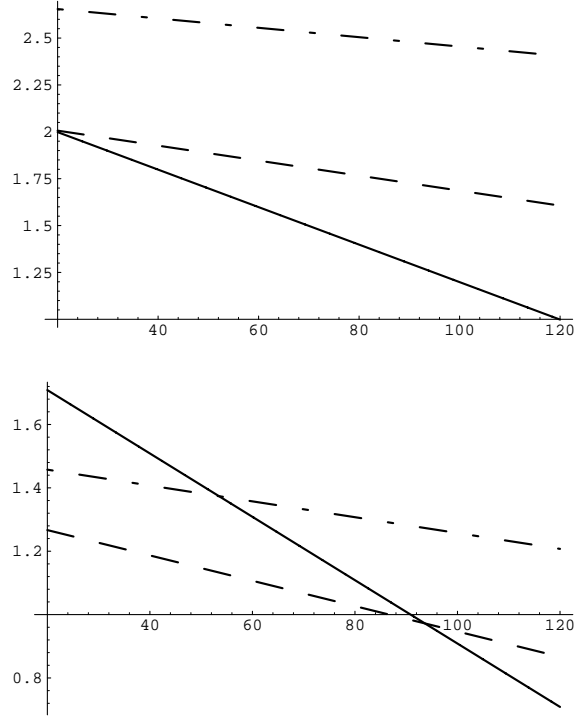


Figure 3: Maximum allowed value of  $A_{\text{eff}}$  in the low  $\mu$  region as a function of  $\mu \cot \beta$  consistent with  $m_{\text{stop}}^{\text{lim}} = 45$  GeV, for  $\tilde{m} = 100$  (solid), 250 (dashed), and 400 GeV (dot-dashed). Upper plot uses the stop  $mass^2$  matrix with equal susy-breaking diagonal terms; lower plot uses modified diagonal entries as discussed in the text.

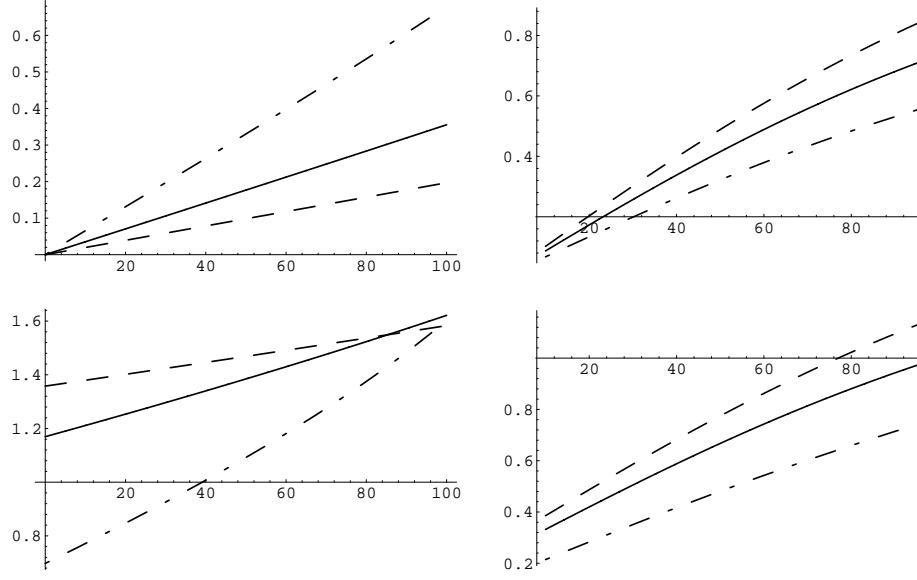


Figure 4: Gluino (left column) and  $\chi_1^0$  (right column) masses in GeV as a function of  $\mu$  in GeV, for  $A_{\text{eff}} = 0$  (upper row) and  $A_{\text{eff}} = 1$  (lower row), with  $\tilde{m} = 100$  (dot-dashed), 250(solid) and 400 GeV(dashed);  $\beta = \frac{\pi}{4}$ .

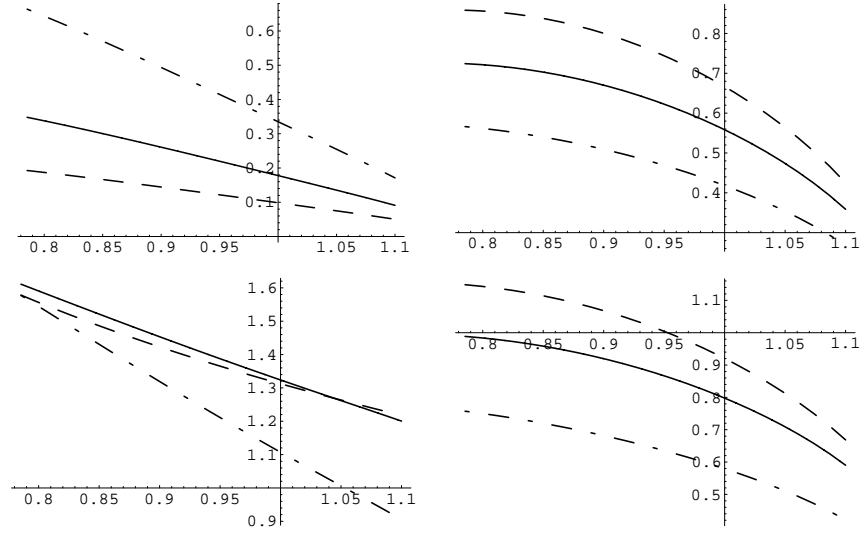


Figure 5: Gluino (left column) and  $\chi_1^0$  (right column) masses in GeV as a function of  $\beta$ , with  $\mu$  taken at its maximum value consistent with the charginos and stops being heavier than 45 GeV. For  $A_{\text{eff}} = 0$  (upper row) and  $A_{\text{eff}} = 1$  (lower row), with  $\tilde{m} = 100$  (dot-dashed), 250(solid) and 400 GeV(dashed).

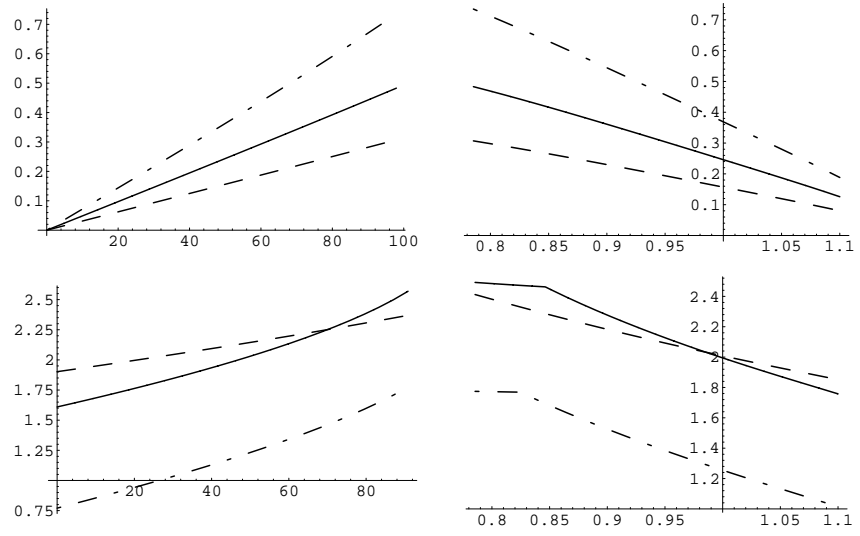


Figure 6: Gluino masses when diagonal SUSY breaking stop  $mass^2$  terms are not taken equal, as described in the text. Left (right) column of plots are as in Fig. 4 (5).

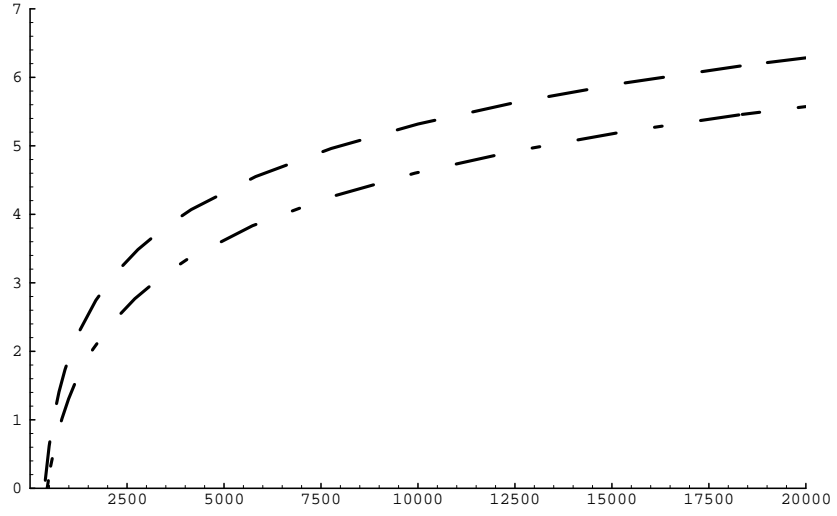


Figure 7: Maximum possible one-loop contribution to the gluino mass in the large- $\mu$  region when  $A_{\text{eff}} = 0$  and  $\beta = \frac{\pi}{4}$ , plotted versus  $\mu$ , corresponding to taking the lighter stop mass to be 45 GeV (dashed) and 126 GeV (dot-dashed) respectively.

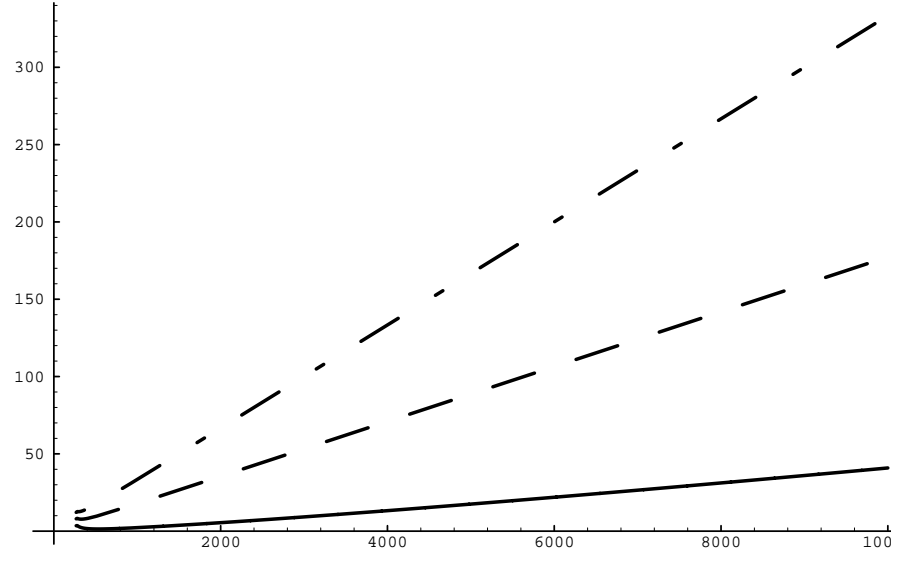


Figure 8: Gluino mass from the two-loop divergent diagram when  $A = 1$ ,  $B = 0$  and  $\beta = \frac{\pi}{4}$  versus  $\tilde{m}$  up to 10 TeV, for  $M_{\text{initial}} = \tilde{m}$  (solid),  $M_{\text{initial}} = 10^{11} \text{ GeV}$  (dashed), and  $M_{\text{initial}} = M_{G1}$  (dot-dashed).



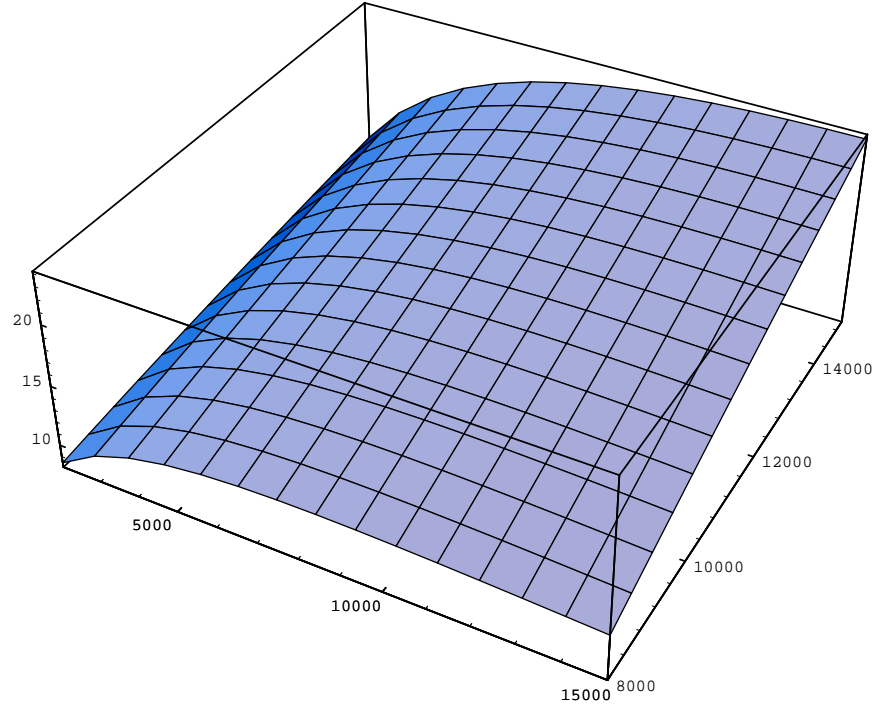


Figure 9: Mass of the lightest neutralino versus  $\mu$  and  $M$  with  $\mu$  running from 2 TeV to 15 TeV and  $M$  from 8 TeV to 15 TeV, taking  $M_{initial} = \tilde{m}$ .